August, 2015

A New Dynamical Picture for the Production and Decay of the XYZ Mesons

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I introduce an entirely new dynamical description for exotic charmoniumlike hadrons, based upon the competing effects of the strong attraction between quarks in a diquark, and the inability of the diquark to hadronize on its own due to being a color nonsinglet. This mechanism naturally explains, for example, the strong preference of the $Z(4475)$ to decay to $\psi(2S)$ rather than the J/ψ, the existence of a state $X(4630)$ that decays to Λ_c baryon pairs, and why some but not all exotics lie near hadronic thresholds. Owing to high-energy constituent counting rules, the four-quark nature of the states produces major changes to both the high-s scaling of cross sections for producing such states and to the potency of the cusp effect of attracting resonances to pair-production thresholds. The recently observed P_c^+ pentaquark candidates are seen to fit naturally into this scheme.

PRESENTED AT

The 7th International Workshop on Charm Physics
(CHARM 2015)
Detroit, MI, 18-22 May, 2015

[1]Work supported by the National Science Foundation under Grant Numbers PHY-1068286 and PHY-1403891.

1 Introduction

In addressing an audience that specializes in charm physics, it is hardly necessary to emphasize the importance of the discovery of "exotic charmoniumlike states" such as $X(3872)$ and $Z(4475)$. Evidence continues to mount that they are tetraquark $c\bar{c}\, q_1 q_2$ states, the first hadrons not classifiable under the $q\bar{q}$ meson/qqq baryon scheme. Approximately 20 exotic charmoniumlike states, at various levels of confirmation, have been observed at BaBar, Belle, BESIII, CDF, CLEO, CMS, D∅, and LHCb (for recent reviews, see Refs. [1, 2]). The situation has become even more interesting with the LHCb observation [3] of pentaquark $c\bar{c}uud$ candidates $P_c^+(4380)$ and $P_c^+(4450)$.

The nonobservation of unambiguous QCD exotics until only a few years ago is one of the more perplexing facts in the history of hadronic physics. QCD allows for the formation of many more color-singlet combinations than just those in the $q\bar{q}/qqq$ paradigm: glueballs, hybrids, tetraquarks, pentaquarks, and so on. Presumably, understanding the way in which these new states are assembled from their quark components will provide key evidence for which phenomena are allowed by QCD and which are not.

In the case of the tetraquarks, a variety of different physical pictures have been proposed. When only the neutral X and Y states were known, it was possible to suppose they were all either conventional $c\bar{c}$ charmonium at unexpected masses or hybrid $c\bar{c}g$ (g = valence gluon) states; but the discovery of the charged Z states mandated a 4-quark structure, at least for those states. However, as time progressed, hadronic transitions between the X, Y, and Z states were observed, suggesting a common structure. Moreover, hybrids are constructed easily only in certain J^{PC} channels such as 1^{++}, while exotics in various J^{PC} channels have now been observed.

The remarkable proximity of some of these states to 2-meson thresholds has spawned three major physical pictures to describe them. Consider, *e.g.*, $X(3872)$:

$$\begin{aligned} m_{X(3872)} - m_{D^{*0}} - m_{D^0} &= -0.11 \pm 0.21\,\mathrm{MeV}\,, \\ m_{X(3872)} - m_{J/\psi} - m_{\rho^0_{\rm peak}} &= -0.49 \pm 0.30\,\mathrm{MeV}\,, \\ m_{X(3872)} - m_{J/\psi} - m_{\omega_{\rm peak}} &= -7.88 \pm 0.21\,\mathrm{MeV}\,. \end{aligned} \tag{1}$$

In comparison, the deuteron, which is considered a loosely bound pn state, has a binding energy of 2.22 MeV, some *twenty times* larger than the central value in the first of Eq. (1). One natural explanation for these numerical oddities is that the tetraquarks are molecules of two (color-singlet) mesons, held together by residual "color van der Waals" forces. In this picture, the binding is accomplished by the exchange of light mesons, particularly pions. This *meson molecule* picture is the most popular of all paradigms for the tetraquarks, having been studied in hundreds of papers. However, the proximity of the charmonium-unflavored meson threshold for many of the exotics also suggests the possibility of *hadrocharmonium* [4], a picture in

which a compact charmonium state lies embedded in a light-quark hadronic cloud, and retains much of its identity until the decay. The values in Eq. (1) also suggest the possibility that at least some of the exotics are not true resonant states, but rather an effect caused by the rapid opening of meson-meson thresholds, which creates a peak in the production rate near the threshold resembling that due to a true resonance; this phenomenon is the so-called *cusp* or *threshold* effect, which has been known for decades in light-quark systems, but was first applied to the charm sector in Ref. [5].

In this talk, however, I wish to advocate for yet one more physical picture for the tetraquarks, which is based upon a well-known yet under-appreciated feature of QCD. The strongest attraction between two color-fundamental quarks is that between a color-$\mathbf{3}$ quark and a color-$\overline{\mathbf{3}}$ antiquark into a color singlet, which is of course the basis of the color structure of a meson. However, it is not the only attractive channel: Two *quarks* can combine into an attractive color-$\overline{\mathbf{3}}$ combination, and the strength of this attraction at short distances is fully half as large as that of the singlet channel. One therefore expects, at least in some physical circumstances, the formation of fairly compact *diquark* states within hadronic systems.

In a system with two quarks and two antiquarks, one therefore has two natural ways to assemble the state: Either one has two associated color-singlet quark-antiquark pairs, as in the molecular or hadrocharmonium pictures, or one pairs the quarks into a diquark, and the antiquarks into an antidiquark. This *diquark picture* for exotic charmoniumlike states was first studied in Ref. [6], and was greatly improved to reflect the results of more recent experiments in Ref. [7]. The greatest difficulties with this picture and the others listed above are summarized in Sec. 2.

Since Refs. [6, 7] discuss the tetraquark states in terms of spin structure in a Hamiltonian formalism, they implicitly treat the tetraquark as a diquark-antidiquark molecule. The specific diquark picture to be described here, introduced in Ref. [8], treats the tetraquark as a type of bound state not previously discussed: The diquark and antidiquark do not actually orbit one another, but remain bound together solely through color confinement. The means by which such states form and decay is discussed in Sec. 3. The use of this picture to probe the multiquark nature of states and its combination with the cusp effect, as well as to provide an explanation of the new P_c^+ states, is briefly discussed in Sec. 4. Conclusions appear in Sec. 5.

2 Limitations of Tetraquark Pictures

Each of the major structural pictures described above for the exotic states presents some dynamical or phenomenological difficulty. Suppose first that the observed states are something other than true tetraquark resonances. In this context, we have already identified the limitations of the hybrid picture. The cusp effect does produce phenomena that resemble resonant peaks, but the narrow $X(3872)$ width ($\Gamma < 1.2$ MeV)

appears too small to be accommodated by a pure cusp not combined with a true resonance [5]. Furthermore, the phase motion measured for the $Z(4475)$ [9] (as well as the P_c^+ states [3]) appears to be consistent with that of a true resonance.

As for true 4-quark bound states, let us begin with a straw man, a simple democratic molecule of the four quarks. In addition to this entity having all the classical instabilities of a 4-body mechanical system, such a molecule would instantly segregate into attractive pairs rather than maintain roughly equal spatial separations [10]. The easy access to these "fall-apart" channels leads one to consider the remaining pictures: hadrocharmonium, meson molecules, and diquarks.

The hadrocharmonium picture was developed to explain the strong coupling of several of the exotics to conventional J/ψ or χ_c charmonium states. However, the couplings of the exotics to $D^{(*)}\overline{D}^{(*)}$ appear to be just as important [indeed, dominant for $X(3872)$]. Moreover, it is unclear why the embedded $c\bar{c}$ pair would remain dynamically stable with respect to the light-quark cloud.

The diquark picture, on the other hand, tends to overpredict the number of bound states, due to its rich color structure; dynamical assumptions (such as which spin couplings dominate [7]) are necessary to reduce the number of states. Moreover, the forces assembling the diquarks, while strong, are still smaller than those between $q\bar{q}$ pairs into color singlets. One would expect diquark molecules with typical interquark separations to re-segregate into meson molecules.

The meson molecular picture is extremely attractive because of its simplicity and the remarkable proximity of many exotics to two-hadron thresholds, as in Eq. (1). However, several of the exotics lie far from such thresholds [*e.g.*, $Z(4475)$], and others lie slightly *above* thresholds [*e.g.*, $Y(4260)$ about 30 MeV above $m_{D_s^*} + m_{\overline{D}_s^*}$], casting doubt on them being bound states. However, the most difficult problem for the meson molecular model is that the most-studied exotic, the $X(3872)$, despite being extremely weakly bound if it is indeed a molecule [again, see Eq. (1)], is produced in large amounts (*prompt production*) in high-energy colliders such as the LHC [11]. In such experiments, quarks are nearly never created with sufficiently small $p_\perp$ to form a state as delicately bound as $X(3872)$; and while final-state interactions (*i.e.*, π exchanges between D^0 and $\overline{D}^{*0}$) substantially expand the $p_\perp$ range that allows molecules to form [12, 13], they do not seem to be enough to explain $X(3872)$ production [14, 15].

3 The Dynamical Diquark Picture

The dynamical diquark picture of Ref. [8] is motivated by several interesting features of the charmoniumlike exotics. As discussed above, the static diquark picture seems problematic due to the possibility of immediate recombination into meson molecules. In order to avoid this recombination, we propose that the diquark-antidiquark (δ-$\bar{\delta}$) pair, once created, rapidly separate to distances r at which the overlaps with meson

wave functions become small. That is, hadronization occurs through the exponentially suppressed large-r tails of the mesonic wave functions, in which δ supplies the quarks and $\overline{\delta}$ the antiquarks. Since each diquark carries color, the δ-$\overline{\delta}$ pair cannot separate indefinitely, but rather convert their kinetic energy into the potential energy of a color flux tube stretching between them.

In Fig. 1 we exhibit the proposed production mechanism for the $Z^-(4475)$ [previously called $Z^-(4430)$] in the process $B^0 \to Z^-(4475)K^+$. One immediately sees that the tetraquark state in this picture is an entirely new type of bound state: not a molecule whose components occupy well-defined orbits, but a dynamical object whose diquark components separate and are distinguishable only due to the large initial kinetic energy imparted to them (here, via $\overline{b} \to \overline{c}c\overline{s}$). One then sees that either path to hadronization, to $(D^{(*)} + \overline{D}^{(*)})$ or to (charmonium + light meson), requires both mesons to have a spatial wave function extent at least as large as the separation between the diquarks. The wave function suppressions lead to a suppressed transition amplitude, and hence a suppressed observable width.

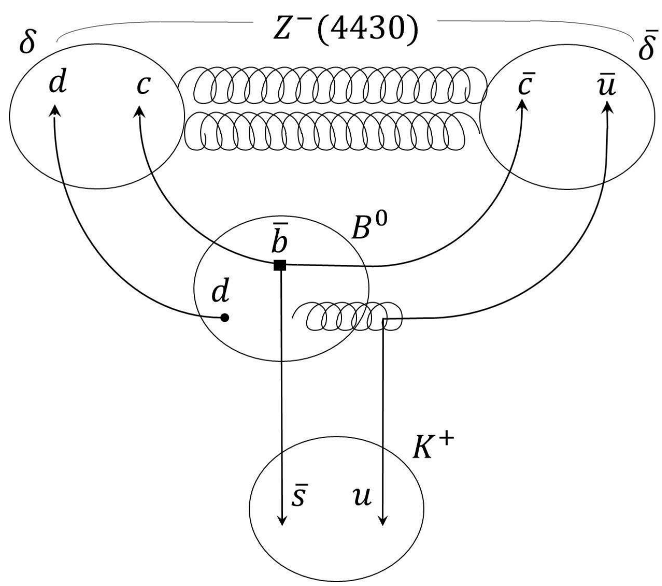

Figure 1: Production mechanism for the $Z^-(4475)$ [from $B^0 \to Z^-(4475)K^+$], where δ and $\overline{\delta}$ indicate the diquarks. The black square indicates the weak decay $\overline{b} \to \overline{c}c\overline{s}$.

Of course, if a B^0 meson decays via $\overline{b} \to \overline{c}c\overline{s}$, the simplest decay modes to construct are two-body modes like $D^{(*)-}D_s^{(*)+}$. Even then, due to the large phase space available, each two-body channel accounts for only about 1%, or collectively for about 5%, of the total decay rate [16]; in the typical B^0 $\overline{c}c\overline{s}$ decay, multiple hadrons are produced. The two-body charmonium decay modes, such as $J/\psi\, K^{*0}$, individually occur at the 10^{-3} level, and collectively around 2.3%. The branching fractions into $X(3872)$

or $Z^-(4475)$, occurring at levels ranging up to $1.7 \cdot 10^{-4}$ for exclusive channels, are at very reasonable levels for states with a δ-$\bar{\delta}$ origin, considering the somewhat smaller attraction within δ or $\bar{\delta}$ compared to that within a color-singlet meson.

Some elements in a picture like Fig. 1 are familiar from textbook discussions of confinement. When the components of a confined state are produced with a large relative momentum (as in jets), they stretch a flux tube between them until the color field contains enough energy to form hadrons by creating an additional $q\bar{q}$ pair from the vacuum (*string fragmentation*). In the picture of Fig. 1, the fragmentation occurs as soon as the threshold for creation of the lightest baryon pair, $\Lambda_c^+ \overline{\Lambda}_c^-$ ($2m_{\Lambda_c} =$ 4573 MeV), is passed. Indeed, the lightest exotic above this threshold, $X(4630)$, is observed to decay predominantly into $\Lambda_c^+ \overline{\Lambda}_c^-$, exactly as expected in this picture.

In principle, hadronization can occur at any point during the δ-$\bar{\delta}$ separation. However, the standard WKB semiclassical approximation predicts the transition probability to be maximal near the classical turning point, *i.e.*, when the kinetic energy of the δ-$\bar{\delta}$ pair converts entirely into the potential energy of the flux tube. The question then becomes how far apart the δ and $\bar{\delta}$ separate before coming to rest. To estimate this distance r, we use that δ and $\bar{\delta}$ are somewhat compact (by virtue of each containing a heavy c quark) color-triplet states and employ the famous linear-plus-Coulomb Cornell potential [17], which has been quite successful [18] in explaining the conventional charmonium spectrum:

$$V(r) = -\frac{4}{3}\frac{\alpha_s}{r} + br + \frac{32\pi\alpha_s}{9m_\delta^2}\left(\frac{\sigma}{\sqrt{\pi}}\right)^3 e^{-\sigma^2 r^2} \mathbf{S}_\delta \cdot \mathbf{S}_{\bar{\delta}} \,, \tag{2}$$

where $\alpha_s = 0.5461$, $b = 0.1425$ GeV^2, $\sigma = 1.0946$ GeV, and $-4/3$ is the color factor specific to $\mathbf{3}$-$\overline{\mathbf{3}}$ attraction. The δ mass can be estimated from lattice or QCD sum rule calculations, but in any case it is close to the D meson mass, and similarly for its charge radius (~ 0.4 fm). Applied to the $Z(4475)$, this procedure gives $r =$ 1.16 fm, and 0.56 fm for $X(3872)$. In comparison, Eq. (2) applied to charmonium gives $\langle r_{J/\psi} \rangle = 0.39$ fm and $\langle r_{\psi(2S)} \rangle = 0.80$ fm. One therefore has a simple explanation of a remarkable experimental fact: The $Z(4475)$ decays to $\psi(2S)$ at least 10 times more frequently than to J/ψ [19], despite both states having the same quantum numbers and much greater phase space available to the J/ψ: It is a simple matter of the large δ-$\bar{\delta}$ state having a greater wave function overlap with $\psi(2S)$ than with J/ψ.

4 Applications

4.1 Dynamical Diquark Resonances and the Cusp Effect

Why should the dynamical diquark picture produce resonant states, and why should these states lie close to the meson-meson thresholds? Three clues noted above lead to

a partial answer to this question: First, the δ-$\overline{\delta}$ states should have observably small widths (*e.g.*, compare $\Gamma[Z(4475)] = 180 \pm 31$ MeV to $\Gamma[\rho(770)] = 150$ MeV). Second, the δ diquarks and D mesons are expected to have similar masses, so the corresponding threshold masses are similar. Third, the cusp effect can combine with existing resonances to drag the observed resonant mass toward meson-meson thresholds [5]. The latter point, especially, is examined in Ref. [20], where it is seen that each such threshold can easily drag a resonant mass several MeV towards it (even, possibly, overshooting the threshold and creating a slightly above-threshold resonance). In the case of the $X(3872)$, Eq. (1) shows that several thresholds are clustered together, potentially creating a strong compound cusp that attracts a resonance to this mass.

Indeed, inasmuch as the diquarks in the δ-$\overline{\delta}$ state can separate substantial distances before being forced to hadronize despite being confined, they approach the behavior of asymptotically free states and can generate a threshold cusp of their own. Since the form factor for creating a multiquark state is expected to fall off at a different rate than that for mesons due to QCD *constituent counting rules* (first discussed in Refs. [21, 22]), the diquark cusp often turns out to be broader and more effective at attracting resonant poles than the meson cusp (see Fig. 2). The full detailed spectrum of the charmonium sector may turn out to be due to a rich combination of "bare" resonances with locations jostled about by both meson-meson and δ-$\overline{\delta}$ thresholds.

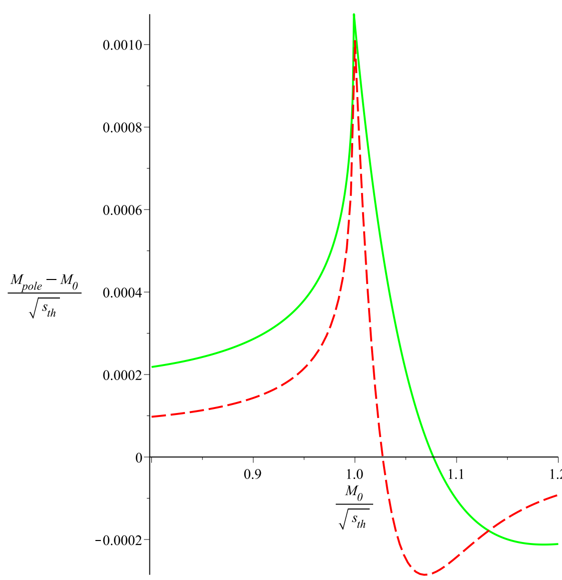

Figure 2: Comparison of the effectiveness of resonant pole dragging by cusps as a function of $M_0/\sqrt{s_{\text{th},i}}$, from the diquark cusp (solid, green) and from the mesonic cusp (dashed, red). Here, $\sqrt{s_{\text{th},i}} = 3.872$ GeV, M_0 and M_{pole} are the bare and final positions of the resonance pole masses, and couplings have been scaled to give the diquark and meson cusp functions the same height.

4.2 Exotics and Constituent Counting Rules

The constituent counting rules are obtained from the twist dimension of the interpolating fields that refer to the hadrons at short distances. They determine the Mandelstam s power-law dependence of cross sections and form factors for processes at large s and fixed scattering angle $\theta_{\rm cm}$: The power of s is determined by the total number n of fundamental constituents (incoming plus outgoing) appearing in the hard scattering. In essence, they amount to counting the number of large-energy propagators necessary to effect the finite-angle scattering of all the constituents. In particular, the invariant amplitude $\mathcal{M}$ for such a process scales as [23]

$$\mathcal{M} \propto \frac{1}{s^{\frac{n}{2}-2}} \, . \tag{3}$$

From Eq. (3), the electromagnetic form factor of a charged tetraquark state such as $Z_c^+ = Z(4475)$ is seen to scale at large s as

$$F_{Z_c^+}(s) \to \frac{1}{s^{\frac{1}{2}(1+1+4+4)-2}} = \frac{1}{s^3} \, . \tag{4}$$

This result is used, *e.g.*, in creating Fig. 2. Here we have taken the natural expectation of 4 fundamental constituents for a tetraquark state. However, if the Z_c^+ contains diquarks that are so tightly bound that they act as fundamental units in high-energy scattering processes, then one expects $F_{Z_c^+}(s) \to 1/s^1$.

While the scaling rules strictly hold only for large s (presumably several GeV above production threshold), their reach may be extended to lower energies by taking the ratios of cross sections of processes that differ primarily through the number of fundamental constituent components, thus eliminating systematic corrections common to both processes. As an example, the ratio

$$\frac{\sigma(e^+e^- \to Z_c^+(\bar{c}c\bar{d}u) + \pi^-(\bar{u}d))}{\sigma(e^+e^- \to \mu^+\mu^-)} = \left|F_{Z_c,\pi}(s)\right|^2 \propto \frac{1}{s^{n-4}} \, , \tag{5}$$

scales as $1/s^4$ if Z_c^+ acts as a two-quark, two-antiquark bound state, while if the diquarks are particularly tightly bound and act as fundamental constituents in the hard scattering, the scaling drops to $1/s^2$. Similarly, consider the ratio

$$\frac{\sigma(e^+e^- \to Z_c^+(\bar{c}c\bar{d}u) + \pi^-(\bar{u}d))}{\sigma(e^+e^- \to \Lambda_c(cud)\overline{\Lambda}_c(\bar{c}\,\bar{u}\bar{d}))} \propto \frac{1}{s^0} \, , \tag{6}$$

such that the same number of constituents, as well as the same heavy-quark ($\bar{c}c$) constituents, appear in both processes. In this case, not only the high-s scaling but also corrections due to the heavy-quark mass cancel in the ratio. One expects the absolute numerical value of the ratio to be substantially smaller if Z_c^+ behaves as a meson-meson molecule than a δ-$\bar{\delta}$ state since the color forces in the former are of the residual van der Waals type and hence much weaker.

4.3 The P_c^+ Pentaquark Candidates

Suppose, with reference to Fig. 1, that one replaces the B^0 meson with the baryon $\overline{\Lambda}_b$, so that the d quark is replaced by the quark pair $\overline{ud}$ [24]. Indeed, let us take the charge conjugate of the diagram so that one may work with baryons rather than antibaryons. The light-quark pair in any Λ_Q baryon ($Q = s, c, b$) has, since the earliest days of QCD, been considered to be a diquark, because it is easily seen to be a "good" spin-0, isoscalar, color-$\overline{\mathbf{3}}$ combination. The presence of the heavy quark b tends to confine the (ud) to a small space (one estimate for the root-mean square matter radius of Λ_b is 0.22 fm [25]), so the diquark $\delta' = (ud)$ can be considered compact, acting essentially as a spectator in the same way as the d in Fig. 1.

The work in Ref. [23] also argued that one can build up more complicated multiquark states (pentaquarks, hexaquarks, *etc.*)—many of which would be easily produced at facilities such as the upgraded JLab—by exploiting the attraction of sequentially formed color-$\mathbf{3}$ and -$\overline{\mathbf{3}}$ combinations. In the case under discussion here, the δ' can combine with a color-$\overline{\mathbf{3}}$ $\bar{c}$ to form a compact color-$\mathbf{3}$ *antitriquark* $\theta = \bar{c}(ud)$. The same reasoning as that in Sec. 3 surrounding Fig. 1 then shows that a *pentaquark* with valence quark structure $c\bar{c}uud$, the result of a rapidly separating color-$\mathbf{3}$ antitriquark $\theta = \bar{c}(ud)$ and color-$\overline{\mathbf{3}}$ diquark $\delta = (cu)$, is an absolutely natural result of the picture. It is the central point of Ref. [24] that this description explains the newly discovered states [3] $P_c^+(4380)$ and $P_c^+(4450)$.

As with the diquark-antidiquark picture for tetraquark states, the new diquark-triquark picture has a static antecedent [26], which was used to explain the then-extant pentaquark candidate $\Theta^+(1535) = u\bar{s}udd$. However, that model also made use of a color-$\mathbf{6}$ diquark inside the triquark in order to achieve a desirable level of binding energy with respect to the nearby KN threshold.

5 Conclusions

We have proposed an entirely new dynamical picture for understanding the exotic charmoniumlike states, such as $X(3872)$, $Z(4475)$, $P_c^+(4380)$, and $P_c^+(4450)$ that have been discovered in recent years and are still being uncovered today. We propose that at least some subset of them are bound, but not molecular, states of color-$\mathbf{3}$ and color-$\overline{\mathbf{3}}$ compact diquarks and triquarks, which have achieved substantial separation due to the large energy release of the process in which they are formed. The states remain bound only due to the confinement of these colored components, and can only decay when color-singlet combinations form through the large-r tails of wave functions of mesons or baryons stretching from one colored component to the other.

Such states can be studied through their multiquark nature using constituent counting rules, as well as their potential ability to create threshold cusps. The next stage of investigation will be to explore the dynamics of the diquark and triquark

formation and the mechanism (going beyond simple quantum-mechanical ideas) by which hadronization across the flux tube is accomplished. The question of why such states occur at some masses and not others, producing the rich spectroscopy already observed, must also be addressed beyond the confines of static Hamiltonian models. Nevertheless, the dynamical picture already gives tantalizing hints of what phenomenology might be possible, in an energy regime that until recently was thought to be a very well-understood sector of particle physics.

ACKNOWLEDGEMENTS

This work would not have been possible without the dedicated efforts of my collaborators, Stan Brodsky, Dae Sung Hwang, and Sam Blitz.

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